

Constraints on Massive Tau Neutrinos and Their Cosmological Implications

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Abstract

We point out the following astrophysical consequences of a tau neutrino with mass in the MeV range; (i) if it has a small electric charge which will allow it to become a cold dark matter of the universe, then present limits on the 511 KeV gamma ray line rule out the possibility that it contributes an Ω_{ν_τ} between .1 to 1 making it unsuitable as a cold dark matter candidate; (ii) if an electrically neutral MeV range ν_τ decays to $\nu_e + \chi$ (where χ is a massless boson) , then its lifetime is bounded by SN1987A observations to be within a window $.05 \text{ sec.} \leq \tau(\nu_\tau) \leq 300 \text{ sec.}/(m_{\nu_\tau} \text{ in MeV})$ a range which may be of interest from the point of view of structure formation. Encouraged by the overlap between the range allowed in (ii) above and that required for a proposed structure formation mechanism by Dodelson et al., we search for simple extensions of the standard model, where such masses and lifetimes may arise for natural values of parameters and show that the already existing singlet majoron and low scale left-right models have this property. We also comment on possible familon models for this decay.

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I. Introduction:

Massive Tau neutrinos with masses in the range $1 \text{ MeV} \leq m_{\nu_\tau} \leq 31 \text{ MeV}$ are allowed by the present laboratory measurements [1] and are not in conflict with cosmological mass density bounds [2] and structure formation provided there is some mechanism (involving either decay or annihilation) that reduces their concentration in the early universe to an appropriate level. There are also interesting allowed ranges of lifetimes of MeV range ν_τ 's which maintain the predictions for primordial light element abundances in agreement with observations[3]. It has recently been suggested that one may be able to use such neutrinos for some positive purposes such as being the cold dark matter of the universe [4] or for helping in understanding the nature of large scale structure in the universe[5]. In this paper, we take a closer look at these suggestions to see if they are consistent with some other astrophysical considerations.

One suggestion is that if the ν_τ is a Dirac particle with mass in the MeV range and has an anomalous electric charge[4] (in the range $3 \times 10^{-6} \leq Q_{\nu_\tau} \leq 3 \times 10^{-5}$) , then its concentration in the early universe can be reduced via the annihilation process $\nu_\tau \bar{\nu}_\tau \rightarrow e^+ e^-$ leaving a freezeout residual density smaller than the cosmological bound. The stable remnant ν_τ 's then could serve as the dark matter on cosmological and galactic scales. Note that the supernova bounds on the Dirac neutrino masses[6] do not apply in this case since the right-handed neutrinos having "strong" electromagnetic interactions get trapped in the supernova after their production.

A second suggestion due to Dodelson, Gyuk and Turner[5] is that a ν_τ in the above mass range which decays as $\nu_\tau \rightarrow \nu_e + \chi$, where χ is a mass less boson (say the majoron or the familon) with lifetime in the range of $0.1 \text{ sec} \leq \tau_{\nu_\tau} \leq 10^3 \text{ sec}$ enhances the fraction of energy in the relativistic components in the universe at the time of its decay. This has the effect of delaying the onset of the perturbation growth and structure formation at the corresponding horizon scales. This in turn allows a better over all fit to a pure cold dark matter model scenario to the fluctuation power spectrum deduced from the study of galactic redshifts and from COBE[7] . An important aspect of this scenario is that the ν_e 's produced in this decay can destroy the otherwise overproduced ^4He (and other light elements) and for a certain range of τ and m_{ν_τ} one can achieve an acceptable big bang nucleosynthesis[3, 8].

The properties of ν_τ invoked in the above two cases were motivated by different reasons but they are none-the-less of physical interest from the point of view of current experiments[9] where there is a possibility that the upper limit on the tau neutrino mass can be reduced to the level of 3-5 MeV. This will of course provide the ultimate test of these proposals. Our goal however is to seek other constraints

on these scenarios in order to test their viability. We find that :

(i) The present limits on the 511 keV gamma rays that result from e^+e^- annihilation of any massive charged dark matter would imply that the contribution of these dark matter particles to Ω will at most be ten percent, thereby making such particles not viable as potential dark matter candidates. These arguments apply not just to tau neutrinos but to any massive minicharged particle which is a potential dark matter candidate.

(ii) We then show that the time profile of the observed neutrino pulses from SN1987A at Kamiokande and IMB implies that for a decay $\nu_\tau \rightarrow \nu_i + \chi$ (where $i = e$ or μ) with a sterile χ , the ν_τ lifetime has a lower bound

$$\tau_{\nu_\tau} \geq 0.05 \text{ sec} \quad (1)$$

Moreover, if in the final state of the decay process, the ν_e mode dominates over the ν_μ mode, then one can also derive an upper bound :

$$\tau(\nu_\tau \rightarrow \nu_e + \chi) \leq 30 - 300 \text{ sec. for } m_{\nu_\tau} = 10 \text{ to } 1 \text{ MeV} \quad (2)$$

It is amusing that there is a considerable overlap between the range defined by our bounds and the one preferred by Dodelson et al. from the structure formation considerations[5]. Encouraged by this, we proceed to note that two existing simple extensions of the standard model (the singlet majoron model[10] and the minimal low scale left-right symmetric model[11]) can naturally provide tau neutrino masses and lifetimes in this range. This should provide additional stimulus for taking this scenario for structure formation seriously and bring new urgency to improving the upper limits on the tau neutrino mass. In particular, one model, the minimal singlet majoron model[10] remains viable in this context only if the electron neutrino mass is no less than an eV or so. This can also be tested in the ongoing neutrinoless double beta decay experiments[12].

II. Minicharged tau neutrino as a potential dark matter candidate:

Before entering into the implications of the minicharged tau neutrino, let us first remind the reader that although unconventional, there is no fundamental reason that can prevent us from entertaining the possibility that the tau neutrino can have an electric charge as large as $\simeq 10^{-5}$ (in units of the positron charge)[13]. We recall that SN1987A observations on the other hand imply that the charge on the electron neutrino cannot be larger than $\simeq 10^{-17}$ [14] . Therefore, any theoretical model that is constructed to lead an anomalous charge for ν_τ should respect this stringent

bound on Q_{ν_e} . One way to achieve this is to gauge an anomaly free quantum number $L_\mu - L_\tau$ in the standard model and have an electric charge formula of the type[15]:

$$Q = I_{3L} + \frac{Y}{2} + \epsilon (L_\mu - L_\tau) \quad (3)$$

The electron neutrino in this model is electrically neutral whereas the ν_μ and ν_τ have equal and opposite charges.

Let us now consider the minicharged tau neutrino as the dark matter of the universe and discuss its implications. If the charge on the tau neutrino is less than 10^{-5} or so, its annihilation cross-section will go out of equilibrium below the nucleosynthesis temperature of one MeV and will not effect the remnant cosmological tau neutrino density. However, the ν_τ 's in the galactic disc is a continuous source of positrons. These positrons of energy $E_{e^+} \simeq 1 - 10 \text{ MeV}$ will be confined by the galactic magnetic fields inside the disc. Their collisions with the disc gas will slow them down on time scales much less than the age of the universe and practically all positrons will then annihilate at rest. In one fourth of these annihilations, namely those proceeding from the singlet spin state, two monochromatic γ - ray lines of energy 511 keV will emerge. We can estimate the flux of these γ -rays locally to be roughly:

$$\Phi_\gamma \approx \left[\frac{1}{4} n_\nu^2 \sigma^{\nu\bar{\nu} \rightarrow e^+e^-} \cdot v \right] R_{eff} \quad (4)$$

R_{eff} is an effective distance parameter related to the radius of the galactic disk and is estimated below. Taking the local ν_τ density to be that of the local dark matter i.e. $400 \text{ MeV}/(cm)^3$, we have

$$n_{\nu_\tau} \simeq \left(\frac{400}{m_{\nu_\tau}/\text{MeV}} \right) (cm^{-3}) \quad (5)$$

The $\nu_\tau \bar{\nu}_\tau \rightarrow e^+e^-$ annihilation cross-section for $2 m_{\nu_\tau}$ sufficiently above the threshold can be written as:

$$\sigma \cdot (v/c) \simeq \frac{2\pi\delta_\tau^2\alpha^2}{4 m_{\nu_\tau}^2} \simeq 3 \times 10^{-36} cm^2 \left(\frac{\delta_\tau}{10^{-5}} \right)^2 \left(\frac{\text{MeV}}{m_{\nu_\tau}} \right)^2 \quad (6)$$

For a cylinder like geometry of the galactic disc with a radius $R \simeq 20$ kilopersecs and thickness $D \approx 1$ kilopersec, the effective radius R_{eff} in eq.(4) can be estimated to

be roughly $\approx (R/4\pi)$, leading to a value for $R_{eff} \approx 1.5$ kpersecs. Substituting back into eq.(4), we get for the 511 keV gamma ray flux generated by this mechanism to be:

$$\Phi_\gamma \simeq 20 \left(\frac{\delta}{10^{-5}} \right)^2 \left(\frac{MeV}{m_{\nu_\tau}} \right)^4 cm^{-2} sec.^{-1} sr^{-1} \quad (7)$$

In ref[4], it was shown that the prediction for Ω_{ν_τ} in this model is given by,

$$\Omega_{\nu_\tau} \approx \frac{1}{10} \left(\frac{\delta}{10^{-5}} \right)^{-2} \left(\frac{m_{\nu_\tau}}{MeV} \right)^2 \quad (8)$$

For $\Omega_{\nu_\tau} \leq 1$ eq.(7) and (8) imply that,

$$\Phi_\gamma \geq 2 \left(\frac{MeV}{m_{\nu_\tau}} \right)^2 cm^{-2} sec.^{-1} sr^{-1} \quad (9)$$

For tau neutrino masses less than 10 MeV, eq.(9) implies a lower bound on the 511 keV photon flux more than $\approx 2 \times 10^{-2} cm^{-2} sec^{-1} sr^{-1}$. The experimental study of gamma rays and the 511 keV line in particular has been an ongoing effort for many years. A fit[16] to the recent GRO[17] data implies a flux of 511 keV line in the disc (in the direction away from the galactic center) at the level of $5 \times 10^{-4} cm^{-2} sec^{-1} sr^{-1}$ and would therefore strongly argue against the minicharged tau neutrino scenario for dark matter.

It has already been noted in [15] that there exist independent theoretical arguments against a tau neutrino charge larger than about 10^{-6} if one uses eq.(3). As already mentioned, eq.(3) implies that the muon neutrino and the tau neutrino must have equal and opposite charges in order to satisfy anomaly cancellations. The existing data on $\nu_\mu e$ scattering[18] can then be used to set an upper limit of about 10^{-9} on the electric charge of the muon neutrino and hence by eq.(3) on δ_{ν_τ} . Moreover, electric charge conservation in muon decay implies that the muon also must have slight anomalous charge equal to δ_{ν_τ} . The precision measurements involving the muon then imply severe restrictions on δ_{ν_τ} . For instance, one gets[15] $|\delta_{\nu_\tau}| \leq 10^{-6}$ from the measured value of $g-2$ of the muon³. These considerations also independently exclude the minicharged tau neutrino scenario for cold dark matter, when it is embedded into the gauge model framework.

³ One could also conceivably improve on this bound by a dedicated experiment where one could look for the displacement of the center of mass of the muonium atom in strong electric field.

III. SN1987A limits on the ν_τ life time:

The cooling of the hot compact cores of the stars following gravitational stellar collapse (eventually leading to a type II supernova) is believed to occur via neutrino emission. The observed neutrino pulses from the SN1987A [19] beautifully confirm this expectation in several respects:

(i). The overall $\bar{\nu}_e$ fluence is consistent with a total collapse energy of $3\text{-}5 \times 10^{53}$ ergs that can be associated with the gravitational binding energy of the star. This energy is released in roughly equal amounts via each of the six neutrino species: ν_e, ν_μ, ν_τ and their anti-particles.

(ii). The energy spectrum of the observed $\bar{\nu}_e$'s is consistent with an expected thermal spectrum with temperature $\approx 4 \text{ MeV}$.

(iii). The duration of the pulse $\Delta t \approx 5\text{-}12 \text{ sec.}$ can be interpreted as the "trapping" time of the initial energetic neutrinos in the dense core. The neutrinos diffuse gradually into the less dense regions and also "cool off" in the process. The effectiveness of the trapping which is governed by $\simeq n_{e,p,n} \times \sigma_\nu$ diminishes as the number density of electrons and nucleons $n_{e,p,n}$ goes down simultaneously with the fact that σ_ν also goes down as the neutrino energy decreases. As a result, the neutrinos escape after a time Δt . The demand that no additional, equally efficient cooling mechanism exists so that the above successful predictions and in particular point (iii) will not be lost, leads to new bounds on interactions that provide new cooling mechanisms.

The two body decay with the sterile massless boson χ (majoron or familon) mentioned earlier provides one such cooling mechanism since the massless boson χ escapes as soon as it is produced⁴. If the decay time is much shorter the above Δt i. e. (denoting $\tau_{\nu_\tau} \equiv \tau$) $\tau \cdot \gamma \equiv \tau(E_{\nu_\tau}/m_{\nu_\tau}) \ll \Delta t \simeq 5 - 12 \text{ sec}$, then this is indeed an efficient cooling mechanism. Demanding that this not be the case yields a lower bound on τ of about .1 sec or so.

Conversely, if ν_τ of mass about 10 MeV decays predominantly into $\nu_e + \chi$ with a relatively long lifetime ($\geq 100 - 1000 \text{ sec.}$), then the decay ν_e would lead to a delayed pulse in the underground detector beyond $\Delta t \simeq 5 - 12 \text{ sec.}$ This reasoning parallels the one made in connection with the 17 keV neutrino [20] yields an upper bound on τ .

⁴In principle, the emitted χ 's can be captured on the ambient ν_e 's and reform ν_τ 's as "resonances". The relevant cross-section at the peak of the resonance is $\sigma_{res} \approx \frac{4\pi}{m_{\nu_\tau}^2}$. However, the probability for having in the $\nu_e \chi$ collision exactly the $E_{cm} = m_{\nu_\tau} \pm \Gamma/2$ with $\Gamma \approx (\text{sec.})^{-1} \simeq 6 \times 10^{-22} \text{ MeV}$ is very small. The back reaction rate is proportional to an effective cross-section $\sigma_{eff} \simeq \sigma_{res} \frac{\Gamma}{E_\nu} \approx 10^{-45} \text{ cm}^2$ for $E_\nu \simeq m_{\nu_\tau} \simeq 10 \text{ MeV}$. This value of σ_{eff} is much smaller than the ordinary weak interaction cross-section that traps neutrinos.

Let us now elaborate on both these arguments . First we discuss the derivation of the lower bound. During $\delta t = \tau \cdot \gamma = \tau(E/m)$, half the energy in the $\nu_\tau + \bar{\nu}_\tau$ components would escape the core due to the decay to χ . This energy can be written as:

$$W_\chi = \frac{1}{2}(W_{\nu_\tau} + W_{\bar{\nu}_\tau}) = \frac{1}{6}f_B W_{tot} \quad (10)$$

In eq.(10), f_B is the Boltzman factor, which is $\simeq 1$ for low masses of the tau neutrino (e. g. $m \simeq 1$ MeV) since the tau neutrino temperature is about 5 MeV. For $m_{\nu_\tau} \simeq 10$ MeV, we can approximate $f_B \simeq (\frac{m}{T})^{\frac{3}{2}}e^{-\frac{m}{T}} \approx .36$. Since the thermal equilibrium population of the tau neutrinos will be immediately regenerated inside the neutrino sphere, the W_{tot} stored in the core will decrease decrease by a factor e over a time

$$t^x(\text{cooling}) = (6/f_B) \times (\tau \cdot E/m) \quad (11)$$

Using $E = 4T_{\nu_\tau} \simeq 20$ MeV, we get $t^x \simeq 120\tau$ for a one Mev tau neutrino whereas $t^x \simeq 40\tau$ for a 10 Mev mass. In order to make this extra cooling mechanism essentially ineffective, we demand that it be bigger than about 12 sec. This leads to the lower bound on τ of .1 sec as mentioned above.

Turning now to the upper bound, we noted earlier[20] that if a ν_τ of mass m decays via a two body mode such as $\nu_e + \chi$ (with both these particles massless) with a lifetime τ , then the ν_e 's arrive earth after a time

$$\delta t = (\frac{m}{2E_{\nu_e}})\tau \quad (12)$$

where E_{ν_e} is the energy of the electron neutrino. Since on the average, $2E_{\nu_e} = E_{\nu_\tau} = 20$ MeV, we have for $m = 1 - 10$ MeV, $\delta t = (.05 - .5)\tau$. These delayed ν_e 's being roughly of similar energy to the original neutrinos will generate a delayed neutrino pulse with roughly 4-13 events in the IMB and 6-19 events in the Kamiokande detector[20]. Since neither detector found any delayed pulse during the time interval of 15 sec. to an hour after the initial pulse, eq.(12) implies an upper bound on τ of about 300 sec. for $m_{\nu_\tau} = 1$ MeV and 30 sec. for $m_{\nu_\tau} = 10$ MeV. As the tau neutrino mass increases, there is suppression of the emitted number of tau neutrinos by the appropriate Boltzman factor and that reduces the overall magnitude of the delayed signal . This for instance implies that the overall quality of our bound for a 10 MeV ν_τ is somewhat weaker than for a 1 MeV ν_τ . Using the estimates carried out in the second paper of ref.[20], we would expect 10-32 delayed events in both experiments for $m_{\nu_\tau} = 1$ MeV as compared to 3.5-11 events for $m_{\nu_\tau} = 10$ MeV.

We wish to point out that the lower bound does not apply if the final state of the tau neutrino decay involves three ν_e 's , since unlike the χ 's, the ν_e 's will get

trapped in the supernova due to their weak interactions and no dramatic cooling can be expected. As for the delayed neutrino pulse, on the average, the final state ν_e 's will have one third of the ν_τ energy and since in the detector the ν_e cross-section goes like $\approx E_{\nu_e}^2$, the secondary delayed pulse will be too weak to yield an upper bound on τ_{ν_τ} .

Finally, in closing this section, we note that the SN1987A favored range for the lifetimes $.1 \text{ sec.} \leq \tau \leq 300 \text{ sec.}$, coincides with that favored by Dodelson et al in their decaying ν_τ CDM scenario. Encouraged by this fortuitous coincidence, we explore in the next section the question of accomodating such lifetimes in simple particle physics models.

IV. Gauge models with desired ν_τ mass and lifetime:

In this section, we consider simple extensions of the standard model which yield ν_τ mass in the MeV range and lifetime in the range of 1 - 100 sec. motivated by the discussions of the previous sections. We will consider see-saw type models which provide the simplest way to understand the smallness of neutrino masses. Ignoring small mixing angles between different generations in the first approximation, the generic mass formula for the light neutrinos in these models has the form:

$$m_{\nu_i} \simeq \left(\frac{m_{\nu_i^D}^2}{f_i v_{BL}} \right) \quad (13)$$

In eq.(13), $m_{\nu_i^D}$ are the Dirac masses of the neutrinos; v_{BL} is the scale of $B-L$ symmetry breaking and f_i are the yukawa coupling of the right-handed neutrinos. In order to estimate the neutrino masses in such models, an additional input about the magnitude of these Dirac masses is needed. Only in grand unified theories, they can be related to the charged fermion masses; but in near electro-weak TeV scale theories that we will be interested in here, they are free parameters. However, motivated by eventual possible grand unification it is conventional to assume that even in TeV scale theories, the neutrino dirac masses are of the same order of magnitude as the charged lepton masses of the corresponding generation. If we then assume that $v_{BL} \simeq 100 \text{ GeV}$ and $f_i \simeq 1$, one gets the following order of magnitude estimates for the different neutrino masses; i.e. ν_τ mass of order $\simeq \text{few MeV}$; ν_μ mass, about a 100 keV and ν_e mass, a few eV. In order to satisfy the cosmological constraints, the model must provide a mechanism for the ν_τ and ν_μ to decay sufficiently fast. We will see below that in the two simplest TeV scale see-saw models, although there are mechanisms for them to decay; the decay rate for ν_μ is much too slow to satisfy the

lifetime bounds coming from structure formation[21] although it can easily satisfy the mass density constraints. Since one of our main motivations for considering this scenario is structure formation, we will take the constraints implied by it seriously. This then implies that the muon neutrino mass must also be in the eV range and for a TeV scale see-saw this requires $m_{\nu_\mu^D} \simeq m_{\nu_e^D} \simeq m_e$. Such a situation could arise if there is some kind of family symmetry between the first and the second generation.

A second point we wish to make concerns the familon models[22] for ν_τ decay. For a global family symmetry that operates on the third generation, present limits on $\tau \rightarrow e + f$ imply a lower limit on the scale of family symmetry breaking, $v_F \geq 10^6 GeV$ or so[23]. If we identify the scale v_F with the $B - L$ symmetry breaking scale v_{BL} , (as one would be tempted to do in simple models), then the see-saw formula produces only a keV range mass for the ν_τ . Therefore, any familon model for an MeV ν_τ must decouple the $B - L$ and the Family symmetry breaking scales.

Let us now to proceed to analyze the ν_τ lifetimes in the singlet majoron and the low scale left-right models.

A. The singlet majoron model:

The singlet majoron model[10] is the most minimal extension of the standard model with nonvanishing neutrino masses and the see-saw mechanism. In addition to the particles of the standard model it consists of three right-handed neutrinos , one for each generation and an additional electroweak singlet Higgs boson Δ^0 that carries total lepton number, $L = 2$. The Lagrangian of the model is assumed to obey the total lepton number symmetry which is spontaneously broken by the Δ^0 field acquiring a vacuum expectation value(vev) at a scale around 100 GeV to a TeV. The neutrino mass matrix then has the see-saw form given in eq.(13) and the earlier discussion about the neutrino spectrum applies. Let us therefore go on to discuss neutrino decays in this model.

The possible relevance of the majoron decay modes of heavy neutrinos for cosmology has been considered in several papers[24]. In particular, using the results of the last paper in ref.[24], we get for $\tau_{\nu\tau}$

$$\tau_{\nu\tau}^{-1} \simeq \sin^2 2\beta \frac{m_{\nu_e} m_{\nu\tau}^4}{16\pi v_{BL}^4} \quad (14)$$

In eq.(14), β is a mixing angle different from the mixing angles measured in the neutrino oscillation experiments. We will choose $\beta = \pi/4$ in order to maximize the decay width in eq. (14). We see from eq. (14) that for a 10 MeV ν_τ , allowing the

ν_e mass of about 2 eV to be consistent with neutrinoless double beta decay bounds, we get $\tau_{\nu_\tau} \simeq 150 \text{ sec.}$, which is at the edge of the allowed range of ref.[5]. As ν_τ mass increases, the lifetime can get shorter and move well into the allowed range depending on the ν_e mass. However due to quartic power dependence on the ν_τ mass, values of m_{ν_τ} smaller than 10 MeV will lead to lifetimes outside the desired range. This is important since it is expected that in the B-factory, tau neutrino masses down to a few MeV can be probed [9]. Similarly, the mass of the ν_e must not be too much smaller than an eV or so if this model is to remain viable. The model is therefore testable in near future.

B. The left-right symmetric model:

Let us now discuss the ν_τ lifetime in the minimal left-right symmetric model with a see-saw mechanism for neutrino masses [11]. Let us remind the reader about the leptonic and Higgs sector of the model. The three generations of lepton fields

are $\Psi_a \equiv \begin{pmatrix} \nu \\ e \end{pmatrix}_a$, where $a = 1, 2, 3$. Under the gauge group $SU(2)_L \times SU(2)_R \times$

$U(1)_{B-L}$, they transform as $\Psi_{aL} \equiv (1/2, 0, -1)$ and $\Psi_{aR} \equiv (0, 1/2, -1)$. The Higgs sector of the model consists of the bi-doublet field $\phi \equiv (1/2, 1/2, 0)$ and triplet Higgs fields: $\Delta_L(1, 0, +2) \oplus \Delta_R(0, 1, +2)$.

The gauge symmetry is spontaneously broken by the vacuum expectation

values: $\langle \Delta_R^0 \rangle = V_R$; , $\langle \Delta_L^0 \rangle = 0$; and $\langle \phi \rangle = \begin{pmatrix} \kappa & 0 \\ 0 & \kappa' \end{pmatrix}$. As usual, $\langle \phi \rangle$

gives masses to the charged fermions and Dirac masses to the neutrinos whereas $\langle \Delta_R^0 \rangle$ leads to the see-saw mechanism for the neutrinos in the standard way and again the discussion about the neutrino spectrum in the introduction to this section applies.

The interactions responsible for neutrino decay in this model comes from the left-handed triplet sector of the theory. As indicated above, these fields do not take part in the Higgs mechanism. The Yukawa Lagrangian relevant to our discussions is given (in the basis where all the leptons are mass eigenstates) by

$$\mathcal{L}_Y = \nu_L^T F' C^{-1} \nu_L \Delta^0 - \sqrt{2} \nu_L^T F'' C^{-1} E_L \Delta_L^+ - E_L^T F C^{-1} E_L \Delta^{++} + h.c \quad , \quad (15)$$

where $\nu = (\nu_e, \nu_\mu, \nu_\tau)$, $E = (e, \mu, \tau)$; F , F' and F'' are 3×3 matrices related to each other as follows:

$$KF = F''; \quad KFK^T = F' \quad , \quad (16)$$

where K is the leptonic CKM matrix in the left-handed sector. First, we note that exchange of the Δ^0 enables the heavier neutrinos to decay to the lighter ones[25]; secondly, the off-diagonal elements of K are the $\nu_e \nu_\mu$, $\nu_e \nu_\tau$ and $\nu_\mu \nu_\tau$ mixing angles, which are directly measurable parameters, and are at present restricted to be, $\theta_{\mu\tau} \leq 0.03$, $\theta_{e\mu} \leq .03$ and $\theta_{e\tau} \leq .17$ [26]. Furthermore, the present upper limits on the $\mu \rightarrow 3 e$ decay can be satisfied by demanding $F_{\mu e} \simeq 0$, since otherwise a large non-zero $\mu \rightarrow 3 e$ decay could arise via Δ_L^{++} exchange. Moreover, in several recent papers, [27] it has been pointed out that there are strong upper limits on the decay mode $\nu_\tau \rightarrow e^+ e^- \nu_e$ from SN1987A observations. In order for the left-right model to be consistent with this bound, we must set $F_{\tau e} = 0$. This leaves only one off-diagonal coupling $F_{\tau\mu}$ free. Present upper limits on the $\tau \rightarrow \mu + \gamma$ decay can be satisfied for $F_{\tau\mu} \simeq 10^{-2}$ or so.

Before discussing the expectations for ν_τ lifetime we wish to note an important property of the model, which is that the decay amplitudes for $\tau \rightarrow \mu ee$ and $\nu_\tau \rightarrow \nu_\mu \nu_e \nu_e$ are not related as can be seen from eq.(15). In fact, even in the limit of vanishing amplitude for $\tau \rightarrow \mu ee$ decay (assured by setting $F_{\tau\mu} = 0$), the ν_τ decay can proceed via the mixing angle $K_{\tau\mu}$. The implication of this observation is that we can obtain a desired value for τ_{ν_τ} without running into conflict with the upper limits on the rare decays of the τ -lepton.

Let us now discuss this in detail. Note that the lifetime for the decay $\nu_\tau \rightarrow \nu_\mu \nu_e \nu_e$ is given by:

$$\tau_{\nu_\tau \rightarrow \nu_\mu \nu_e \nu_e}^{-1} = \frac{F_{\tau\mu}'^2 F_{ee}^2 m_{\nu_\tau}^5}{768 \pi^3 M_{\Delta^0}^4} \quad (17)$$

If we choose $F_{\tau\mu} \simeq 0$ in order to forbid the rare μee decay of the tau lepton, then in eq. (17) $F_{\tau\mu}' \simeq \theta_{\tau\mu} \times (F_{\tau\tau} - F_{\mu\mu})$. Using the present limits on the neutrino mixing angles, we can have $F_{\tau\mu}'$ of order 0.06. Assuming $F_{ee} \simeq .1 - 1$, we then get tau neutrino lifetime of about .2 to 20 sec. which is in the range desired in [5]. Only if the upper limit on $\theta_{\tau\mu}$ goes below 10^{-3} , then one has to invoke a non-zero $F_{\tau\mu}$ (and hence a nonzero branching ratio for $\tau \rightarrow \mu ee$) to get the desired ν_τ lifetime of order 100 sec.

Another point worth noting is that unlike the singlet majoron model, the lifetime of ν_τ in this case is independent of the ν_e mass. Therefore, this model will remain viable even if the upper limits on the ν_e mass go down to the level of .1 eV from future neutrinoless double beta decay searches.

V. Conclusion:

We have pointed out two astrophysical constraints on an MeV range tau neutrino: the first one has to do with an unconventional MeV ν_τ which carries a small electric charge that could make a potential dark matter candidate. Our investigation shows that this interesting possibility runs into conflict with the data on 511 keV gamma rays obtained with the help of the Gamma Ray Observatory. The second one deals with the conventional electrically neutral ν_τ with an MeV range mass, which is unstable. We show that if decays to $\nu_e + \chi$, then SN1987A neutrino observations imply an upper as well as lower bound on its lifetime. This allowed window is interesting in the sense that it coincides with one considered desirable by Dodelson et al.[5] in their attempts to understand structure in the universe using only cold dark matter plus a decaying MeV range tau neutrino. We then discuss several gauge models where such lifetimes can arise naturally in the hope that this new scenario for structure formation will receive more serious consideration by both cosmologists as well as particle physicists.

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